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Coronal outflow dominated accretion discs: a new possibility for low luminosity black holes?

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ABSTRACT
The spectral energy distributions of galactic black holes in the low/hard state and of low-luminosity AGN possess many common features, the most prominent being: compact, flat (or inverted) spectrum radio cores with high brightness temperatures; excess red and infrared emission, often correlated with the radio flux; an extremely weak (or absent) quasi-thermal hump and a hard X-ray power-law with high energy cut-off. These sources are thought to be accreting at low rates and advection (or convection) dominated accretion flows are usually considered the best candidates to explain them. Here we present an alternative possibility, involving strong, unbound, magnetic coronae generated by geometrically thin, optically thick accretion discs at low accretion rates. First we show that, if angular momentum transport in the disc is due to magnetic turbulent stresses, the magnetic energy density and effective viscous stresses inside the disc are proportional to the geometric mean of the total (gas plus radiation) and gas pressure. Therefore the corona is less powerful in a radiation pressure dominated disc, and the relative fraction of the power liberated in the corona increases as the accretion rate decreases. Furthermore, we discuss reasons why energetically dominant coronae are ideal sites for launching powerful jets/outflows, both MHD and thermally driven. In analysing the spectral properties of such coronal outflow dominated accretion discs, we reach the important conclusion that if the jet/outflow is, as is likely, radiatively inefficient, then so is the source overall, even without advection of energy into the black hole being relevant for the dynamics of the accretion flow.

Key words: accretion, accretion discs – black hole physics – magnetic fields

1 INTRODUCTION: ADAF OR JETS ?

The spectral energy distribution (SED) of the gravitational power released as electromagnetic radiation when matter accretes onto a compact object is far from universal. Different accretion modes are possible, and often the same initial conditions at the outer boundaries admit more than one solution for the accretion flow configuration at the inner boundary, with often very different radiative properties. The main goal of accretion flows theory is to distinguish and understand all the possible different modes of accretion, and classify the observed sources in terms of such modes.

In the case of non-quiescent Galactic (stellar mass) Black Holes Candidates (GBHC), an X-ray based classification in terms of their *spectral states* has emerged in recent years. Broadly speaking, the different states are distinguished by their broadband luminosity, and by the relative contribution to the luminosity of a 'soft', quasi-thermal and a 'hard', power-law like spectral components (see e.g. the recent review by Done 2001 and reference therein). At luminosities close to the Eddington one, these sources are often in the *Very High State*, where both of the two components contribute substantially to the SED. At slightly lower luminosities the quasi-thermal one dominates, and the power-law is usually steeper and extended to the gamma-ray band (*High/Soft State*). At even lower luminosities the spectra are completely dominated by a hard power-law component, with the quasi-thermal one extremely weak or even absent: these are the so-called *Low/Hard States*. Sometimes, at luminosities intermediate between those of the Soft and the Hard States, an *Intermediate State* is observed, with properties similar to those of the Very High State. It is often the case that the same source, either persistent or transient, undergoes a transition between spectral states, and therefore accretion modes.

From the theoretical point of view, the quasi-thermal component is usually associated with a geometrically thin, optically thick accretion disc. Its dynamical and geometrical properties could in principle be deduced solely by a careful comparison between the observed spectral properties of the quasi-thermal hump and a detailed model for the accretion disc. On the other hand, the observed hard X-ray power-laws represent a universal signature of a specific physical process (most likely inverse Compton scattering of soft photons on a population of hot electrons; Sunyaev & Titarchuk 1980) more than of specific accretion dynamics. This is the reason why, if there is little doubt that the standard accretion disc model (Shakura & Sunyaev 1973) captures the basic physical properties of black holes in their High/Soft state, the accretion mode responsible for the Low/Hard state is still matter of debate.

Recently, multi-wavelength observations of GBHC in the

Low/Hard state have been adding important pieces of information, which are helping to shed more light on the subject. Crucial in this respect is the discovery of the high-galactic latitude transient source XTE J1118+480. This source was persistently in the Low/Hard state during its 2000 outburst, and being very little affected by galactic extinction, has provided the most accurate and complete SED to date of a galactic black hole in this state (McClintock et al. 2001; Frontera et al. 2001). It possesses all the typical characteristics that distinguish such state (Zdziarski 2000; Fender 2001), namely: a hard X-ray power-law with photon index $\Gamma \simeq 1.8$ and a high energy cut-off, indicative of thermal Comptonization; a dim and cold thermal disc component peaking in the EUV band (McClintock et al. 2001); excess red and infrared emission compared the one expected from standard disc models; a compact (unresolved) flat (or inverted) radio core, probably extending to the NIR or optical regime (Fender et al. 2001).

The situation for supermassive black holes accreting in the centre of active galactic nuclei bears some resemblance with the GBHC one. Beside the best studied case of the black hole at the centre of our Galaxy (Melia & Falcke 2001, and references therein), recent multi-wavelength campaigns on nearby low-luminosity AGNs (LLAGN) (Ho 1999, Nagar et al. 2000, Falcke et al. 2000, Ho et al. 2001) have revealed nuclear SEDs markedly different from the canonical broadband continuum spectrum of high luminosity AGN (as compiled, e. g., by Elvis et al. 1994). The most striking difference is the lack of UV excess, the so-called 'big blue bump' associated with the emission from an optically thick, geometrically thin accretion disc (Ho 1999). In addition, VLA (Nagar et al. 2000) and VLBA (Falcke et al. 2000) radio observations of LLAGN have revealed flat or inverted compact radio cores, with an incidence higher than in the case of normal Seyfert galaxies, and, for the sources observed at higher spatial resolution with VLBA, very high brightness temperatures ($T_b \gtrsim 10^8$K).

In light of the above observational similarities between the two classes, we would like to discuss here a possible answer to the following question: what is the common accretion mode for low-luminosity black holes?

Since their rediscovery in recent years (Narayan & Yi 1994), radiatively inefficient, optically thin, advection dominated accretion flows (ADAF; Ichimaru 1977, Rees et al. 1982, Abramowicz et al. 1995; see Narayan, Mahadevan & Quataert 1999 for a recent review) have been regarded as natural solutions. They can explain the absence of the quasi-thermal hump, the low X-ray luminosities and the observed X-ray spectral slopes. When tested against the best data for GBHC in their Low/Hard state, though, as in the case of XTE J1118+480 (Esin et al. 2001) or Cygnus X-1 (Esin et al. 1998), ADAF models alone cannot work. A transition between an inner ADAF and an outer Shakura Sunyaev disc is needed, with $R_{\rm tr} \approx 10 - 100 R_{\rm S}$, as can also be inferred from studies of X-ray reflection components (Gilfanov, Churazov & Revnivtsev 2000; Done 2001, and reference therein). Detailed spectral fits to the broadband SED of the nearby LLAGNs M81 and NGC 4579 (Quataert et al. 1999) lead to similar conclusions. The physics of such a transition, which may be taking place through a gradual evaporation of the disc as proposed by Meyer & Meyer-Hofmeister (1994) or via turbulent diffusive heat transport in the radial direction (Honma 1996), is not included in the ADAF models, so that $R_{\rm tr}$ is usually treated as a free parameter in the spectral fit procedure. Furthermore, the low energy part of the SED is also difficult to reconcile with the standard ADAF picture, that in general predicts highly inverted spectra, instead of the observed flat ones. In fact, morphology and spectral indices of the compact radio cores require additional outflows from the central inner part in order to be consistent with the prediction of an adiabatic flow (Di Matteo, Carilli & Fabian 2001). Finally, we must note that, both from the theoretical point of view (Abramowicz, Lasota & Igumenshchev 2000) and from numerical simulation (Igumenshchev, Abramowicz, & Narayan 2000), it is clear that radiatively inefficient flows are subject to strong convective instabilities (and might therefore be better called CDAF; Stone, Pringle & Begelman 1999, Quataert & Gruzinov 2000a), but are not in general able to generate strong outflows (although, see Misra & Taam 2001), unless the viscosity is very high.

Flat or slightly inverted radio spectra from very compact sources are generally associated with jet models (Blandford & Königl 1979; Reynolds 1982). This has prompted the alternative view that spectra of LLAGN are essentially jet-dominated (Falcke 2001). As a matter of fact, in low-luminosity nearby radio galaxies observed with HST (Capetti et al. 2000) compact optical cores seem very common. Chiaberge et al. (1999), analysing HST images of a complete sample of 33 FR I radio galaxies, demonstrated that the luminosity of compact optical cores in these sources correlates linearly with that of the radio cores, suggesting a common non-thermal synchrotron emission origin for both. In particular, detailed spectral jet models have been demonstrated to successfully reproduce the SED of Sgr A* (Falcke & Markoff 2000). Interestingly, this might also be the case for the X-ray transient XTE J1118+480, (Markoff, Falcke & Fender 2001). The presence of a relativistic jet in the black hole candidate Cygnus X-1 in the low/hard state has been recently confirmed by milliarcsecond resolution VLBA observations (Stirling et al. 2001), while Fender (2001), on the basis of the observational evidence of strong radio compact jet correlating with hard X-ray emission in GBHC in the Low/Hard state, concluded that, in these sources, the jet power is at least comparable with the accretion one.

Here we explore further this latter hypothesis from the point of view of the standard accretion disc theory. We shall propose a possible new model for low-luminosity black holes, in which a standard geometrically thin, optically thick disc at low accretion rates dissipates a large fraction of its gravitational energy in a magnetic corona. We shall outline the reasons why energetically dominant coronae are the ideal site for launching powerful jets/outflows, and discuss the expected spectral properties of such accretion flows. The important conclusion we reach is that if the jet/outflow is radiatively inefficient, then so is the source overall.

2 CORONAL STRENGTH AT LOW ACCRETION RATES

It is well known that micro-physical viscosity is in general too inefficient to drive accretion at rates comparable to the observed ones, and enhanced/turbulent viscosity is therefore a theoretical necessity. As astrophysical discs are bound to be seeded with magnetic fields, magneto-rotational instability (MRI) (Velikhov 1959; Chandrasekhar 1960; Balbus & Hawley 1991) is the best candidate to date to produce the self-sustained turbulence that is able to transport angular momentum in rotationally supported discs. Any weak, sub-thermal magnetic field, coupled with outwardly decreasing differential rotation (two conditions easily met by any accretion flow endowed with angular momentum) generate the turbulence which in turn is able to regenerate the field (Tout & Pringle 1992). This instability, which has a rapid growth rate of the order of the orbital frequency Ω, results in a greatly enhanced effective viscosity that

is able to transport angular momentum outward (Balbus & Hawley 1998).

What is the fate of the magnetic field which is amplified inside the disc? Two competing mechanisms should operate in an accretion disc to saturate the field: dissipation and buoyancy. On one hand, part of the field *must* be dissipated, if the turbulence has to be sustained and MRI relied upon to transport angular momentum. In fact it has been shown (Balbus & Hawley 1998) that, in a steady state disc, if angular momentum transport is due solely to magneto-rotational instability, magnetic fields must be sub-thermal. This in turn implies that buoyancy cannot be a too efficient mechanism to get rid of the disc field: to cross one disc scaleheight in less than a dynamical time $1/\Omega$ a velocity of the order of the sound speed is required, which would correspond to equipartition magnetic field. Such strong fields would suppress the MRI (Balbus & Hawley 1998).

On the other hand, we can expect the magnetic field configuration in the disc to be highly intermittent (Politano & Poquet 1995), with the field concentrated in small regions (flux tubes and ropes) with relatively low filling factor (see e.g. Blackman 1996 and the simulations presented in Machida, Hayashi & Matsumoto 2000). These structures can more easily rise buoyantly and emerge in the low density part of the flow above the accretion disc to form a magnetic atmosphere, the corona, where the field is finally dissipated.

It is not clear at present what is the relative importance of the two mechanisms of field saturation inside a standard accretion disc, if they either coexist or instead alternate in a cyclic fashion (as, for example, in the sequence: low disc field → MRI → stronger field → buoyancy → field expulsion → suppression of the MRI → low disc field → MRI; see also the discussion in Tout & Pringle, 1992).

Numerical simulations of vertically stratified discs (Brandenburg et al. 1995; Stone et al. 1996) seem to suggest that, although dissipation is the main saturation mechanism for the instability, buoyancy is indeed capable of transporting vertically a substantial fraction of magnetic energy to power a strong corona (see in particular Miller & Stone 2000, where a larger vertical extent is explored), without quenching the magneto-rotational instability in the underlying disc.

In this section we investigate how changes in the internal disc structure, associated with variation in the external accretion rate, may influence the accretion disc–corona system. To this end we only consider standard optically thick and geometrically thin accretion discs sandwiched by a patchy magnetic corona (Galeev, Rosner & Vaiana 1979; Haardt, Maraschi & Ghisellini 1994; Di Matteo, Celotti & Fabian 1999; Merloni & Fabian 2001b). We will assume that buoyant vertical transport of magnetic flux tubes is a relatively efficient way of disposing of the disc field, and discuss the self-consistency of such assumption at the end of the section.

The nature of the saturation mechanism determines the level up to which magnetic fields are amplified inside the disc. Many authors have argued that magnetic field amplification is likely to be limited to values such that the magnetic pressure is at most equal to the local *gas* pressure, even in its inner, radiation pressure-dominated parts (Sakimoto & Coroniti 1981; Stella & Rosner 1984; Sakimoto & Coroniti 1989; Vishniac 1995a). In a recent paper Blaes & Socrates (2001) have analysed local dynamical instabilities in magnetized, radiation pressure-supported (and geometrically thin) accretion discs, and found that, due to radiative diffusion, the growth rate of the MRI is reduced in such systems to $\sigma \sim \Omega(c_{\rm g}/v_{\rm A})$, where $c_{\rm g}$ is the gas sound speed and $v_{\rm A}$ is the Alfvén speed, well below the usual rapid growth rate ($\sim \Omega$) in gas pressure dominated discs. We can therefore give an order of magnitude estimate of the amplitude of the field in the radiation pressure dominated part of a standard accretion disc. If we assume that the buoyant vertical speed of a magnetic flux tube is approximately given by the Alfvén speed, we get for the buoyant timescale $t_{\rm b} \sim H/v_{\rm A} \simeq c_{\rm s}/(\Omega v_{\rm A})$, $c_{\rm s} = \sqrt{P_{\rm tot}/\rho}$ is the isothermal sound speed. The magnetic field, with initial (subthermal) amplitude B_0, rapidly grows to $B \sim B_0 e^{\sigma t_{\rm b}}$. Therefore we have $\ln(B/B_0) \simeq c_{\rm s} c_{\rm g}/v_{\rm A}^2$. Neglecting the logarithmic dependence on the initial field, this in turn suggests that the field will saturate at an amplitude such that

$$\frac{B^2}{8\pi} = P_{\rm mag} \simeq \alpha_0 \sqrt{P_{\rm gas} P_{\rm tot}}, \tag{1}$$

where $P_{\rm tot}$ and $P_{\rm gas}$ are the total (gas plus radiation) and the gas pressure at the disc midplane, respectively, and α_0 is a constant of the order of unity. Such a scaling was first suggested by Taam & Lin (1984), and its stability properties have been studied by Lin & Shields (1986) and, for the case of slim discs at high accretion rates, by Szuszkiewicz (1990).

If magnetic turbulence is ultimately responsible for transporting angular momentum, we can still retain the usual modified α viscosity prescription in which the stress tensor is proportional to the gas pressure $\alpha = P_{\rm mag}/P_{\rm gas}$ (Sakimoto & Coroniti 1981; Stella & Rosner 1984), but now with $\alpha \simeq \alpha_0 \sqrt{P_{tot}/P_{\rm gas}}$.

As in many other previous treatments of the two-phase thermal models for accretion flows onto compact objects (Haardt & Maraschi 1991; Svensson & Zdziarski 1994), we parameterize the coronal dominance by means of the fraction f of gravitational power associated with the angular momentum transport that is stored in the form of magnetic structures (loops, flux tubes), transported vertically through the disc in a dissipationless fashion, and finally dissipated in the corona. Thus, if the total released power of the accretion disc–corona system is given by $L \equiv \dot{m} L_{\rm Edd} = 4\pi G M m_{\rm p} \dot{m} c/\sigma_{\rm T}$, we define the corona luminosity as $L_{\rm c} = f \dot{m} L_{\rm Edd}$. The fraction f of power dissipated in the corona can be expressed as the ratio of the magnetic Poynting flux in the vertical direction to the locally dissipated energy flux in the disc:

$$f = \frac{F_{\rm P}}{Q} = \frac{P_{\rm mag} v_{\rm P}}{Q}, \tag{2}$$

where $F_{\rm P}$ is the vertical Poynting flux and Q is the dissipation rate of gravitational energy per unit surface area, and we have defined the velocity $v_{\rm P}$ with which the magnetic flux is transported in the vertical direction. From the standard equations of geometrically thin accretion discs with modified viscosity law we have

$$Q = \frac{3}{2} \alpha v_{\rm K} \left(\frac{H}{R}\right) P_{\rm gas} = \frac{3}{2} \alpha c_{\rm s} P_{\rm gas}, \tag{3}$$

where $v_{\rm K}$ is the Keplerian velocity.

The main source of uncertainties lies in the complex process of vertical transport of magnetic flux tubes in the body of the accretion disc, and therefore on the value of the velocity $v_{\rm P}$ (see Coroniti 1981; Stella & Rosner 1984; Vishniac 1995a,b; for a more thorough discussion). Here we assume that the rise speed of the magnetic tubes is proportional to their internal Alfvén speed, and consequently scales as $v_{\rm P}/c_{\rm s} = b\beta^{-1/2}$, where $\beta = P_{\rm tot}/P_{\rm mag}$ is the standard plasma parameter, and b is related to the efficiency of buoyant transport of magnetic structure in the vertical direction inside the disc, which is of the order of unity for extremely evacuated tubes.

Equation (2) then becomes

$$f \simeq \frac{2b}{3}\beta^{-1/2} = \frac{2b\sqrt{\alpha_0}}{3}\left(1+\frac{P_{\rm rad}}{P_{\rm gas}}\right)^{-1/4}. \qquad (4)$$

The fraction f can then be calculated for any given accretion rate and at any given radius. If the mass of the central black hole is $M = mM_{\odot}$ and the radial distance from the central source is expressed in units of Schwarzschild radii $r = R/R_{\rm s} = Rc^2/2GM$, from (4) and the equations for the ratio $P_{\rm rad}/P_{\rm gas}$ with modified viscosity law, we obtain the following implicit equation for $f(r)$ as a function of r in the radiation and gas pressure dominated zones of the disc, respectively (cfr. Stella & Rosner 1984, appendix A)[1]

$$\frac{Kf(r)^{-4}-1}{(1-f(r))^{9/5}} \simeq 2.01\times10^5(\alpha_0 m)^{1/5}r^{-21/10}[\dot{m}J(r)]^{8/5} \qquad (5)$$

$$\frac{Kf(r)^{-4}-1}{(1-f(r))^{9/10}} \simeq 1.31\times10^2(\alpha_0 m)^{1/10}r^{-21/20}[\dot{m}J(r)]^{4/5}, \qquad (6)$$

where the function $J(r) = 1-\sqrt{3/r}$ comes from the no-torque inner boundary condition. The numerical factor $K = (2b\sqrt{\alpha_0}/3)^4$ contains the unknown quantities in our treatment. As we shall see, the value of K is related to the maximal fraction of power released in the corona. It is evident that high viscosities and vertical rise speeds (that are associated with the presence of very low-β, strongly evacuated, magnetic filaments in the disc) are needed in order for the corona to be powerful.

The fraction of power dissipated in the corona is higher when gas pressure dominates over radiation pressure inside the disc. This happens over a larger and larger portion of the disc as the accretion rate decreases. The radial boundary between the two solution is found when the fraction of power released in the corona is $f(r) = (K/2)^{1/4}$, and can be calculated by solving the equation

$$\frac{r}{J(r)^{16/21}} \simeq 350(\alpha_0 m)^{2/21}\dot{m}^{16/21}(1-(K/2)^{1/4})^{6/7}. \qquad (7)$$

The left hand side has a minimum at $r_* = 5.72$, therefore, for accretion rates smaller than a critical value, $\dot{m}_{\rm cr} \simeq 0.016(\alpha m)^{-1/8}(1-(K/2)^{1/4})^{-9/8}$, the disc is gas pressure dominated even in its inner part.

Equations (5) and (6) can be solved numerically to calculate the disc structure at every radius and for every accretion rate. A global value of $\langle f\rangle$ can then be obtained by integrating over all the disc area:

$$\langle f\rangle = \frac{\int_3^\infty f(r)Q(r)2\pi r dr}{\int_3^\infty Q(r)2\pi r dr}. \qquad (8)$$

In Figure 1 we plot such global value of $\langle f\rangle$ as a function of the total accretion rate for two different values of K (0.99 and 0.1) for typical stellar mass and supermassive black holes.

As already pointed out, as the radiation pressure dominated part of the disc shrinks, the value of $\langle f\rangle$ increases as a larger portion of the disc is gas pressure dominated. For accretion rates smaller than the critical one, $\dot{m}_{\rm cr}$, when the disc is completely gas pressure dominated, the $\langle f\rangle$–$\dot{m}$ relation flattens, and the fraction of the power released in the corona increases only slightly as the disc becomes denser and the radiation pressure further decreases. The maximum coronal dominance is attained for the lowest values of the accretion rate, and is roughly equal to $K^{1/4}$. If $K > 1$ eqs. (5) and (6) do not admit a solution with $f(r) < 1$ over a wide range of radii. This might correspond to a situation in which the magnetic field is so strong and/or the vertical transport of magnetic flux tubes is so efficient that the standard accretion disc – corona configuration in no longer stable and the accretion disc, which cannot maintain its Keplerian profile, is effectively fragmented by the magnetic field into a number of disconnected rings (Heyvaerts & Priest 1989). Also, very low-β discs are qualitatively different in terms of angular momentum transport processes. In fact, in such cases the magneto-rotational instability is suppressed and magnetic stresses at the disc–corona interface may be the primary angular momentum vehicle (Blandford & Payne 1982) and the standard disc structure we have assumed may not be appropriate. We therefore did not investigate such a case further, restricting ourselves to the stable case in which $\beta \gtrsim 1$ in the disc, and $K \lesssim 1$.

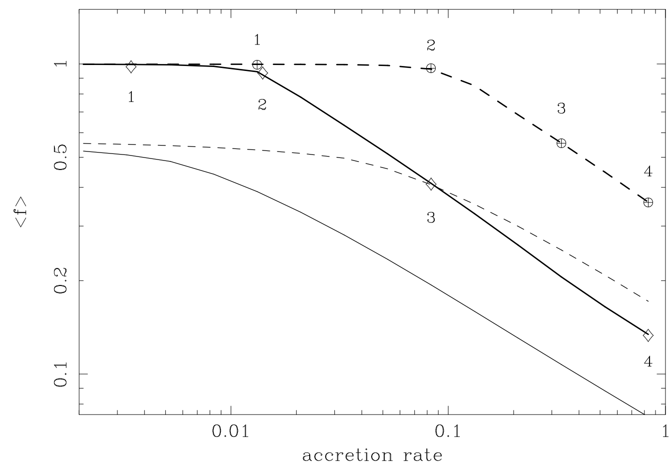

Figure 1. The integral fraction of accretion power dissipated into the corona as a function of the accretion rate for $\alpha_0 = 0.3$ and two different values of central black hole mass: $M = 10^8 M_{\odot}$ (solid lines) and $M = 10 M_{\odot}$ (dashed lines), and of the coefficient K: thicker lines corresponds to $K = 0.99$ and thinner ones to $K = 0.1$. The maximal coronal strength is reached for the smallest accretion rates, and is roughly equal to $K^{1/2}$. The marked numbers on the $K = 0.99$ curves correspond to the spectral energy distributions shown in Figures 5 and 6.

The coronal strength depends weakly on the central source mass, as implicitly shown in eq. (5) and (6). In Fig. 2 we show explicitly the values of the accretion rate for which $f > 0.5$ ($\dot{m}_{5,\rm f} \propto M_{\rm BH}^{-0.124}$). Ghisellini & Celotti (2001b), by studying the dividing line between FR I and FR II radio-galaxies (which correspond approximately to low and high radio–power galaxies, respectively) in the radio–host galaxy luminosity plane conclude that such dividing line correspond to an approximately universal (independent on the black hole mass) transition value of $\dot{m} \simeq 0.06$, consistent with the values we found for the accretion rates at which the corona starts to be the dominant repository of gravitational energy.

To summarize, in accretion discs where angular momentum is transported either by magnetic turbulence, or stresses, magnetic coronae should be dominant at low accretion rates, and their strength depends upon the nature of magnetic dissipation inside the disc. Furthermore, magnetic coronae are stronger in stellar mass black holes than in AGN (see also Fig. 2 below).

[1] We have derived the structure equations for a cold disc on the basis of Svensson & Zdziarski 1994, eqs. 1-4. The parameter ξ of the radiative diffusion equation has been set equal to 1 for the radiation pressure dominated solution and to 8/3 in the gas pressure one in order for the two solutions to match continuously at the boundary. Also, for the ease of computation, we have approximated $(\alpha/\alpha_0)^{1/5} \simeq 1$ and $(\alpha/\alpha_0)^{1/10} \simeq 1$ in the expression for $P_{\rm rad}/P_{\rm gas}$.

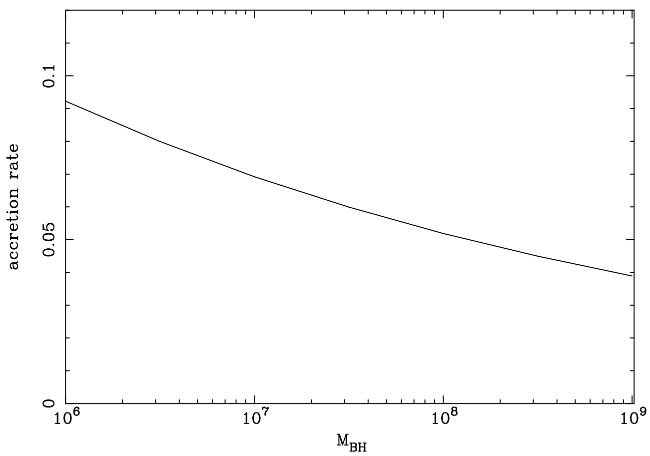

Figure 2. As a function of the central black hole mass, we plot the value of the accretion rate for which the dissipation in the corona is more than the internal disc one ($f > 0.5$).

3 THE FATE OF THE CORONAL ENERGY: CORONAL HEATING AND JET LAUNCHING MECHANISMS

What is the fate of the energy deposited in the corona by the magnetic structures that emerge from the disc when f is large?[2] If the current paradigm for high energy emission from black holes is correct, involving inverse Compton scattering of soft photons on a population of thermal electrons at temperatures $T_{\rm e} \sim 10^9$K, then part of this energy must be used to heat the corona, most likely via magnetic reconnection (Di Matteo 1998) and then efficiently radiated away. On the other hand, because Compton cooling is so efficient in the strong radiation field of an accretion disc, the energy densities of the thermal electrons and of the hard X-ray radiation are likely to be only a small fraction of the total energy stored in the corona, whose dominant contribution should come from stored magnetic field (Merloni & Fabian 2001a) and possibly hot ions, if the corona is two-temperature (Di Matteo, Blackman & Fabian 1997). In such a situation, the condition can be favourable for the occurrence of strong coronal outflows[3]. The nature of the outflow depends crucially on the magnetic field strength and topology and on the physical properties of the particle heating mechanisms.

In Merloni & Fabian (2001a) it has been shown how the values of coronal temperatures and optical depths inferred from the observations of low luminosity black holes require that the main reservoirs of energy in the corona are either a structured magnetic field or the protons in a two-temperature plasma. The extent to which the protons may act as repository of most of the dissipated energy crucially depends on how efficient the heating of electrons in the corona is. Great theoretical effort has been devoted recently to the study of particle energization in accretion flows, mainly for its relevance for the ADAF scenario. Quataert (1998) and Quataert & Gruzinov (1999) have shown that in optically thin, advection dominated accretion flows, as long as the MHD turbulence is dissipated on small scales via incompressible modes (Alfvén waves), protons will be mainly heated only in high β plasmas ($\beta \gtrsim 10$), while electrons are more likely to be heated in low β plasmas. On the other hand Blackman (1999) argued that if modes of dissipation involving compression dominate, then, even for $\beta \sim O(1)$, electron heating is negligible. Furthermore, we stress that the field configuration in the corona ought to be more complicated than the uniform field considered by Quataert (1998) and Quataert & Gruzinov (1999), with the field concentrated in loops and tubes, similarly to those observed in the sun. Magnetic reconnection may play an important role in this case (see e.g. Bisnovatyi-Kogan & Lovelace 1997; Di Matteo 1998). Unfortunately, the physics of particle energization in magnetic reconnection is still poorly understood, and no clear conclusions can be safely drawn at present. However, the conclusion that a powerful corona is the ideal launching platform for powerful jets remains unscathed: as a matter of fact, magnetically dominated coronae ($\beta \lesssim 1$) may power MHD driven outflows which can carry away a substantial fraction of the gravitational power (see section 3.1), even if the ions are subvirial. On the other hand, for $\beta \gtrsim 1$, ions could be the main repository of the power dissipated in the corona; they can be heated to supervirial temperature and launch powerful thermally-driven outflows (see section 3.2).

For the sake of simplicity, we keep the treatment of these two cases separate, bearing in mind that in a realistic situation they may both contribute to the determination of the total outflowing power. We demonstrate that the power carried away by the outflow/jet may be a dominant fraction of the total accretion power in both cases.

To be fairly general, we consider that the total coronal power generated by the accretion disc, $L_{\rm c} = f\dot{m}L_{\rm Edd}$, can be either dissipated locally to heat the corona, and ultimately radiated away as hard X-ray radiation with a luminosity $L_{\rm H} = (1-\eta)L_{\rm c}$, or used to launch a jet with power $L_{\rm j} = \eta L_{\rm c}$. In the following, we investigate the dependence of the parameter η on the disc parameters.

The coronal magnetic field intensity depends on the dissipation rate and on the field geometry. Here we assume, on the basis of the strong spectral and variability arguments (Haardt, Maraschi & Ghisellini 1994; Di Matteo, Celotti & Fabian 1999; Merloni & Fabian 2001b), that the corona is structured and the magnetic field concentrated in a number of small active regions (where dissipation occurs). The size, height and distribution of such regions determines the details of the spectral and temporal properties of the high energy emission in these sources (Merloni & Fabian 2001b). Here it will suffice to notice that a relatively small number ($N \sim 100$, see Merloni & Fabian 2001b) of spherical active regions of size and height of the order of a few Schwarzschild radii ($R_a \sim H_a \sim 1R_{\rm S}$) is what is needed to reproduce the average properties of the hard X-rays emission of GBHC in their hard state and of Seyfert 1s and low luminosity AGN, namely a spectral index $\Gamma \sim 2$, high energy cut-off around 100 keV and strong variability.

We therefore assume that the magnetic energy density in the corona is

$$\frac{B^2}{8\pi} = \frac{3L_{\rm H}}{4\pi R_a^2 N c}\left(\frac{c}{v_{\rm diss}}\right), \tag{9}$$

where $v_{\rm diss}$ is the dissipation velocity and depends on the uncertain nature of the reconnection process (Di Matteo 1998; Merloni & Fabian 2001b), and can be assumed to be of the order $v_{\rm diss} \sim 0.01c$.

3.1 The jet power I. MHD launching mechanisms

Here we discuss the most general mechanisms by which a coronal magnetic field can power strong outflows (either in the form

[2] For the sake of clarity, in the rest of this work we will use f to indicate the global value of the fraction of gravitational power released in the corona $\langle f \rangle$, neglecting any radial dependence.

[3] As we do not discuss either the collimation mechanism or the final velocities of the outflowing gas, but only the power channeled in such a jet/outflow, we do not make any distinction between jets and outflows, and use the two terms indifferently.

of Poynting flux or in the form of a magnetically driven material wind).

Models and simulations of jet production (Blandford & Znajek 1977; Blandford & Payne 1982; Meier 1999) show that it is the *poloidal* component of the magnetic field which mainly drives the production of powerful jets, and the output power[4] can be expressed as (Livio, Ogilvie & Pringle 1999)

$$L_{\rm j} = \left(\frac{B_{\rm p}^2}{8\pi}\right) 2\pi R_{\rm cor}^2 R_{\rm cor}\Omega. \qquad (10)$$

Here $R_{\rm cor}$ is the size of the region where most of the coronal power is dissipated and Ω is the typical angular velocity of the magnetic field lines. Maximal rotation of the central black hole may be important in this context more for its deeper potential well and the larger rotational velocities that the accreting gas can attain (see also Meier 2001) than for the extraction of rotational energy from the black hole itself.

The magnetic field that emerges from the disc into the corona must be mainly azimuthal, as simulations also suggest (Miller & Stone 2000; Machida, Hyashi & Matsumoto 2000). If the disc is not threaded by any external large scale poloidal field, as it is likely to be in the case of central black holes, the strength of the poloidal field component depends on the typical scaleheight of a coronal magnetic flux tube and on the capability of reconnection events to create larger and larger coherent structures. Tout & Pringle (1996) have proposed numerical and analytic models in which stochastic reconnection creates magnetic loops on all scales, with a power-law distribution. Here, as the jet-launching region is identified with the corona itself, we consider the constraints put on the poloidal field by the coronal geometry rather than the disc one.

A magnetic flux tube emerging from the disc should have a cross section of the order of the disc thickness (Galeev, Rosner & Vaiana 1979). Nonetheless, the strong differential rotation at the tube footpoints causes a twisting of the tube magnetic field and increase its tension, counterbalancing its sideways expansion (see Parker (1979), §9.1, 9.6). As both theoretical arguments and numerical simulations have shown (Aly 1990a; Aly 1990b; Amari et al. 1996; Romanova et al. 1998), the net effect of this energy injection in the magnetic coronal loops is therefore a rapid loop expansion, followed by reconnection (at a height H_a much larger than the disc thickness) and opening of the field lines. Then, if H_a is the typical coronal flux tube scaleheight (height of a reconnection site), we have

$$\frac{B_{\rm p}}{B} \simeq \frac{H_a}{R_{\rm cor}}. \qquad (11)$$

From eqs. (9), (10) and (11) we obtain

$$L_{\rm j} = \frac{3}{2} L_{\rm H} \left(\frac{c}{v_{\rm diss}}\right) \left(\frac{H_a}{R_{\rm cor}}\right)^2 \left(\frac{R_{\rm cor}\Omega}{c}\right), \qquad (12)$$

which in turn gives, for the fraction of coronal power that goes into the MHD jet/outflow

$$\eta_{\rm MHD} = \left(1 + \frac{2}{3}\left(\frac{v_{\rm diss}}{c}\right)\left(\frac{R_{\rm cor}}{H_a}\right)^2\left(\frac{c}{R_{\rm cor}\Omega}\right)\right)^{-1}. \qquad (13)$$

Therefore, from eq. (13), we can conclude that the jet/outflow power is stronger if:

(a) The dissipation speed is low. In this case magnetic field does not dissipate effectively via reconnection and can reach high intensities, favouring MHD outflow launching;
(b) The coronal scaleheight is large with respect to the distance from the central source. This would help in increasing the relative strength of the poloidal component of the magnetic field, that is the one ultimately responsible for the powering of the jet;
(c) The rotational velocity of the field lines is high. This is related to the gravitational potential of the central black hole (maximally rotating Kerr holes can increase the angular velocity of field lines relative to a distant observer due to the strong frame dragging effect, see e.g Meier 2001), and to the typical radius of maximal coronal dissipation.

As an example, for $v_{\rm diss} \simeq 0.01c$, $R_{\rm cor} \sim 7R_{\rm S}$, $H_a \sim 2R_{\rm S}$ and $\Omega = \Omega_{\rm K}(R_{\rm cor})$, we obtain $\eta_{\rm MHD} \simeq 0.55$: the MHD jet can carry away a substantial fraction of the coronal (and accretion) power. The fraction of jet power $\eta_{\rm MHD}$ strongly depends on the typical height of the coronal active regions. Interestingly, for a given active region size, R_a, the greater its distance from the disc, H_a, the more important will be synchrotron radiation as a source of soft photons for Comptonization (Di Matteo, Celotti & Fabian 1999; Merloni & Fabian 2001b) as opposed to external radiation from the disc. This immediately suggests that sources with observed variable high frequency synchrotron radiation are strong candidates for harbouring powerful MHD driven jets.

3.2 The jet power II. Thermally driven outflows

As already discussed at the beginning of section 3, a crucial uncertainty about the energetics of powerful coronae is whether dissipation of magnetic energy preferentially heats the protons or the electrons. It seems very likely that the fraction of energy dissipated into the electrons is negligible at least for optically thin plasmas where magnetic pressure is smaller than gas pressure (Quataert & Gruzinov 1999). In this case, as discussed in Di Matteo, Blackman & Fabian (1997), a two-temperature equilibrium (with $T_{\rm p} \gg T_{\rm e}$) is easily established.

Let us therefore consider the case of a two-temperature corona (see also Janiuk & Czerny 2000; Różańska & Czerny 2000) and assume that all the dissipated magnetic energy is dumped into the protons

$$\left(\frac{d\epsilon_B}{dt}\right)_- = \frac{B^2}{8\pi}\frac{v_{\rm diss}}{R_a} = \left(\frac{d\epsilon_i}{dt}\right). \qquad (14)$$

The protons exchange energy with the electrons only via Coulomb collision, so that, in a steady state,

$$\left(\frac{d\epsilon_{\rm i}}{dt}\right) = \left(\frac{d\epsilon_{\rm e}}{dt}\right)_+, \qquad (15)$$

where $(\frac{d\epsilon_{\rm e}}{dt})_+$ is the rate of energy transfer from ions to electrons due to Coulomb collisions between populations with Maxwellian distributions, calculated using the Rutherford scattering cross section (Stepney & Guilbert 1983):

$$\left(\frac{d\epsilon_{\rm e}}{dt}\right)_+ = \frac{3}{2}\frac{m_{\rm e}}{m_{\rm p}}\frac{n_e n_i \sigma_{\rm T}(kT_{\rm i} - kT_{\rm e})c}{K_2(1/\theta_{\rm e})K_2(1/\theta_{\rm i})} \ln\Lambda f(\theta_{\rm e},\theta_{\rm i}), \qquad (16)$$

[4] Based on the conclusion of the analysis presented in Livio, Ogilvie & Pringle (1999), we will neglect the contribution to the total power output from the energy extracted from the black hole via the Blandford–Znajek effect (BZ, Blandford & Znajek 1977). If the angular momentum of the central black hole is high enough, and a large scale poloidal field is generated and sustained in the ergosphere, the BZ power could contribute substantially to the final jet/outflow power, but the main conclusion of our investigation, namely the dependence of the jet power on the disc accretion rate, would remain unchanged (see e. g. Meier 2001).

where $T_{\rm e,i}$ are the electrons and ions temperatures, respectively, $\theta_j = kT_j/m_jc^2$ ($j=$ e, i) are dimensionless temperatures, $\ln\Lambda \approx 20$ is the Coulomb logarithm, Ks are the modified Bessel functions and

$$f(\theta_{\rm e},\theta_{\rm i}) = \left[\frac{2(\theta_{\rm e}+\theta_{\rm i})^2+1}{\theta_{\rm e}+\theta_{\rm i}}K_1\left(\frac{\theta_{\rm e}+\theta_{\rm i}}{\theta_{\rm e}\theta_{\rm i}}\right)+2K_0\left(\frac{\theta_{\rm e}+\theta_{\rm i}}{\theta_{\rm e}\theta_{\rm i}}\right)\right].$$

In order to calculate the electron temperature, we equate the electron heating rate with the total cooling rate $q_{\rm cool} = q_{\rm Comp} + q_{\rm synch} + q_{\rm CS}$, which includes inverse Compton scattering of both disc photons and self-absorbed synchrotron emission produced in the active region itself (see e.g. Di Matteo, Celotti & Fabian 1997, Wardziński & Zdziarski 2000).

Once the heating rate is specified via eqs. (9) and (14), and the geometry of the corona is fixed, we can calculate electron and ion temperatures if we know the coronal optical depth. Observationally, both black hole candidates in their Low/Hard state and AGN have similar values of τ, lying in a narrow range around $\tau \sim 1$ (Gierliński et al. 1997; Zdziarski 1999). As an illustrative example, we made here the assumption that, for each value of the central source mass, M, the coronal optical depth is proportional to the accretion disc surface temperature. In particular, for gas pressure dominated solution appropriate for low accretion rates, we have chosen $\tau \simeq 4\dot m^{2/5}(1-f)^{1/5}$. Therefore τ is smaller for smaller accretion rates ($\tau \sim 0.2$ for $\dot m = 0.005$) and we have $\tau \sim 1$ for $\dot m = 0.05 - 0.08$, depending on the black hole mass. The dependence of τ on the accretion rate for central black holes of masses $m = 10$ and $m = 10^8$ is shown in Fig. 3. It is important to stress that the specific dependence of τ on the disc parameters we have assumed here has been chosen only for illustrative purposes: the results we present on the strength of thermally driven outflows would be qualitatively the same for any physical situation in which colder and denser accretion discs were sandwiched by lower optical depth coronae (as would be the case, for example, if the coronal optical depth depended on the accretion disc ionization parameter $\xi = L_{\rm H}/nH_a^2$). In fact, the only requirement for the jet/outflow to dominate the energy budget at increasingly lower accretion rates is that the coronal optical depth decreases with decreasing accretion rate, regardless of the specific profile. If instead τ was approximately independent of $\dot m$, the highest proton temperature would be reached for $\dot m_{\rm cr}$, where the absolute coronal power, $L_{\rm c} = f\dot m L_{\rm Edd}$ is maximum.

Once a value for the coronal optical depth is fixed, for each value of the accretion rate we can then calculate magnetic field intensity, electron and proton temperatures and the overall emitted spectrum by solving the equation for the cooling and heating rate, adopting an analytic approximation for the emissivities described in Merloni & Fabian (2001b; Appendix A). In Figure 3 we also show, as functions of the accretion rate, the calculated values of the electron temperature $\Theta = kT_{\rm e}/m_{\rm e}c^2$ and spectral index α of the hard X-rays produced by inverse Compton scattering on the hot coronal electrons.

We stress that the values of the electron temperature depends also on the details of the coronal geometry, namely, on the size, height and number of magnetic regions reconnecting in the corona at any time (Stern et al. 1995; Di Matteo, Celotti & Fabian 1999). This is an inherently time dependent problem (see e.g. Poutanen & Fabian 1999, Merloni & Fabian 2001b) and here we have chosen average values for the coronal geometric parameters ($H_a = 2R_a = 2R_{\rm S}$) that ensure that the value of the spectral index is close to the observed one, and we are not interested in detailed spectral modeling of the X-ray emission.

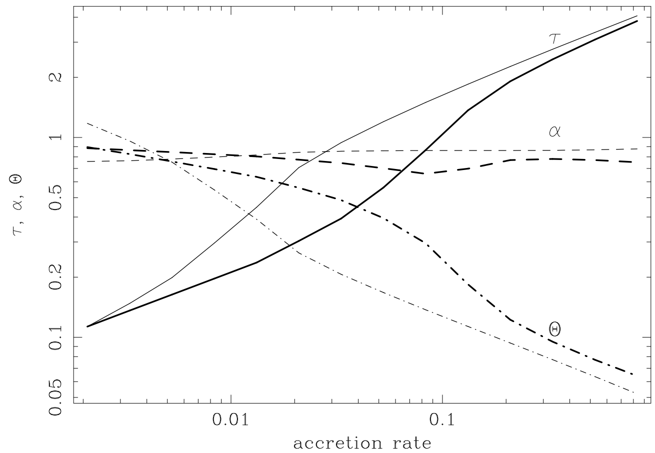

Figure 3. Variation of the electron temperature in the corona ($\Theta = kT_{\rm e}/m_{\rm e}c^2$, dot-dashed lines) and of the X-ray spectral index (α, dashed lines) as functions of the accretion rate, calculated for the optical depth profile shown (τ, solid lines). Thicker lines correspond to a central black hole mass $m = 10$, thinner lines to $m = 10^8$.

If the coronal density is low enough, and if the power dissipated in the corona high (and both conditions are more likely in low accretion rate discs, as we have shown in section 2), the protons can be heated up to temperatures higher than the virial one $T_{\rm vir} = GMm_{\rm p}/3kR_*$ (here we may take as R_* the radius at which the coronal dissipation is higher). In this case the corona is likely to drive powerful jets/outflows (Piran 1977; Lin, Misra & Taam 2001). A fraction η of the coronal power is advected away in the outflow in the form of bulk kinetic energy and the heating rate of the protons is therefore reduced. In such cases we may estimate a value for the parameter η simply by looking for the minimal value of η for which the protons are at the virial temperature. The resulting values for such $\eta_{\rm thermal}$ as a function of the accretion rate for $K = 0.99$ and two different black hole masses are shown in Figure 4. The outflow/jet power (kinetic energy plus Poynting flux) decreases sharply as the accretion rate increases and is negligible as soon as the temperature up to which the protons are heated in the corona drops below the virial one.

4 PREDICTED SPECTRA

Regardless of the nature of the launching mechanism, a robust conclusion from the analysis of the previous sections is that when the fraction of gravitational power released into a structured magnetic corona approaches unity, strong outflows ought to be produced. Here we describe the basic spectral properties of the disc–corona emission. As a reference, we will consider a sequence of models taken from the two–temperature solutions illustrated in Fig. 1 and 4, but we emphasize that any of such spectra would be observed if the jet was MHD driven, as far as it had the same value of η.

For small and intermediate values of f (therefore for $\dot m \gtrsim 0.08$ or $\dot m \gtrsim 0.008$ for $m = 10$ and $m = 10^8$, respectively) the seed soft photons for Comptonization come from the quasi-thermal accretion disc emission (either intrinsic or reprocessed). Then, the coronal cooling decreases when the accretion rate decreases, and the X-ray spectrum gets harder, as can be seen if Fig. 3. The hardest spectra are obtained for $f \to 1$ and $\eta < 0.8$. At this point inverse Compton scattering on synchrotron photons produced in the corona dominates the cooling. Finally, as the ions become supervirial, more and

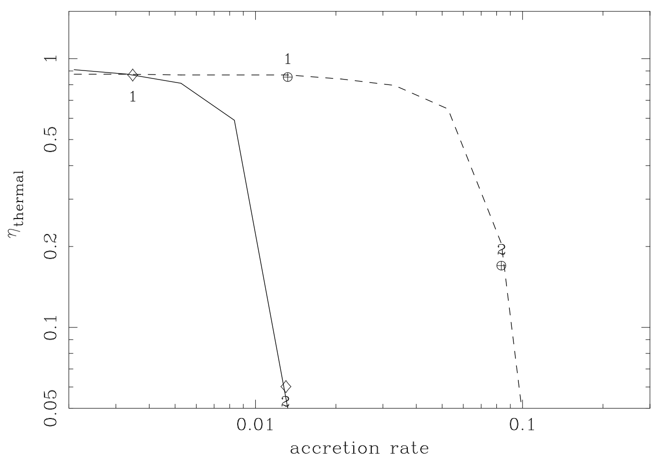

Figure 4. The relative strength of thermally driven coronal outflow as a function of the accretion rate for two different values of the central black hole mass: $M = 10^8 M_{\odot}$ (solid line) and $M = 10 M_{\odot}$ (dashed line). We have imposed maximal coronal strength by taking $K = 0.99$. At the lowest accretion rates, where $f \rightarrow K^{1/4}$ we have $\eta_{\rm thermal} \simeq 0.98$.

more power is carried away in the form of a powerful jet/outflow and the coronal heating is reduced more rapidly than the cooling due to synchrotron emission and the spectrum gets slightly softer again.

This is illustrated in Fig. 5 and 6 which show a sequence of spectral energy distributions for the stellar mass and supermassive black hole cases, respectively.

We do not attempt to model in detail the emission from the jet/outflow. Given the high brightness temperatures measured in the cores of some LLAGN, we can expect self-synchrotron Compton emission to be important at low $\dot{m}$. Also, the high energy spectrum would be modified by even a small fraction of non-thermal electrons in the corona and/or in the outflow (see e. g. Poutanen & Coppi 1998; Wardziński & Zdziarski 2001).

5 DISCUSSION

At the low accretion rates we are considering, both our corona dominated solution for a standard thin disc and ADAF/CDAF solution are admitted. A 'strong ADAF principle' has been advocated (Narayan & Yi 1995) to ensure that the real systems will tend to chose the ADAF/CDAF branch whenever possible. Here we have explored the possibility that such a strong ADAF principle does not hold, and studied the expected properties (spectral and energetic) of the alternative accretion mode. Nonetheless, there are admittedly some similarities between the powerful outflow-dominated, radiatively-inefficient coronae that we found here at the lowest accretion rates and a classical ADAF, and we might well be exploring a region of the accretion flow parameter space where the two accretion modes are bridged together (in this case mainly by virtue of strong magnetic field generation, rather than evaporation; Meyer & Meyer-Hofmeister 1994). Indeed, predictions from ADAF models of a transition in spectral properties at $\dot{m} \lesssim 0.1$ (Esin, Mc Clintock & Narayan 1997) are similar to those of our coronal outflow dominated sources, at least at high energies. Also, from our study we expect strong coronal outflows at high values of the viscosity, similar to what is predicted by the theory of advective flows (Igumenshchev & Abramowicz 2000). The main testable difference between the two models is the presence of a geometrically thick, cold,

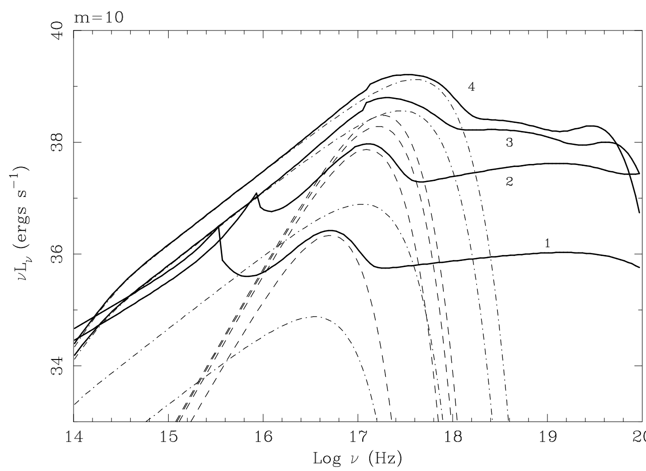

Figure 5. The spectral energy distributions calculated for our thermally driven jet model for $m = 10$ (see text for details). The four spectra correspond to the four points marked with open circles in Fig. 1. Solid lines show the total emerging spectra, while dashed lines show the contribution of thermal reprocessed radiation from the disc and dot-dashed lines the intrinsic multicolor disc emission (with outer radius $R_{\rm out} = 2 \times 10^5 R_{\rm S}$. The strong peak at $10^{15} - 10^{16}$ Hz represent the self absorbed synchrotron emission from the active coronal regions: for spectrum 1 ($\dot{m} \simeq 0.01$) it clearly represents the main source of seed photons for Comptonization. The total jet luminosity has been estimated assuming a flat spectrum ($L_{\nu} \propto \nu^{\gamma}$ with $\gamma \sim 0$) extending up to the frequency of coronal synchrotron self-absorption and a 10 per cent radiative efficiency. No optically thin emission from the jet is included. Also neglected is the reduction of the thermal reprocessed hump (dashed lines) that can be expected if the X-ray emitting coronal active regions also move with the outflow at relativistic speeds (Beloborodov 1999). This is by far the dominant contribution to the EUV/soft X-ray emission at low accretion rates.

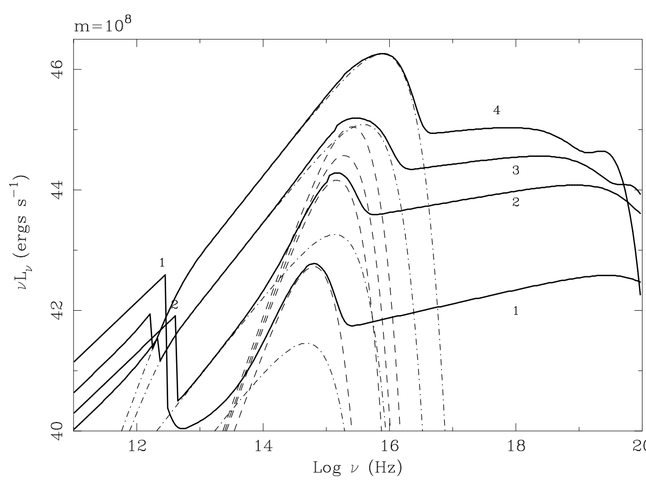

Figure 6. The spectral energy distributions calculated for our thermally driven jet model for $m = 10^8$ (see text for details). Symbols and line styles are the same as in Figure 5. Here direct synchrotron radiation is always dominated by the quasi-thermal disc component, due to the decreased value of the coronal magnetic field ($B \propto m^{-1/2}$), except for spectrum 1, at the lowest accretion rate, where a strong jet component is evident up to about 1000 GHz. At accretion rates lower than what shown by curve 1 ($\dot{m} < 0.003$) the shape of the SED does not change any longer, the overall luminosity being simply reduced.

accretion disc in the inner part of the flow. In the magnetic corona scenario, even at the lowest accretion rates, when most of the accreting power is dissipated in the corona, the presence of the disc in the inner part of the flow would manifest itself by reprocessing coronal radiation and producing both a cold, quasi-thermal and a reflection components. Nonetheless, it is well known that if the coronal luminosity is so high that the inner part of the accretion disc is fully ionized, the reflection features are so weak that it is difficult to distinguish between a truncated disc and an ionized one (Done & Nayakshin 2001; Ballantyne, Ross & Fabian 2001). Furthermore, when the corona is so powerful that it produces relativistic outflows, the bulk motion of the coronal outflowing material reduces the amount of high energy radiation impinging on the disc, thus reducing the strength of the reprocessed components (Beloborodov 1999; Malzac, Beloborodov & Poutanen 2000).

When $f \to 1$ and $\eta \to 1$, the radiative efficiency of a coronal outflow dominated flow depends on the radiative efficiency of the outflow itself. However, unless $1 - \eta \sim 10^{-3}$, we cannot achieve the extremely low efficiencies ($\epsilon \lesssim 10^{-5}$) required to explain some nearby LLAGN (Di Matteo et al. 2001). Such estimates, though, are based on the Bondi accretion rates estimated at relatively large distances from the source. It is plausible to conceive that the strong outflows we envisage, by interacting with the intergalactic medium, would stifle accretion through the inner part of the disc (an effect analogous to that of the ADIOS solution of Blandford & Begelman 1999, in the advective case), reducing the value of the accretion rate through the inner disc that we are considering here and alleviating the need for too low a radiative efficiency (see also discussions in Melia and Falcke, 2001, with regards to the Galactic center, and Di Matteo et al., 2001, for NGC 6166).

Our coronal-outflow-dominated solutions are both thermally and viscously stable, as in general are all standard Shakura–Sunyaev accretion disc solution in the gas pressure dominated regime. Rapid and dramatic variability in the observed high-energy flux is nevertheless expected, as X-rays are produced by coronal structures that are the eventual outcome of the turbulent magnetic field generation inside the disc. The geometry of these structures (open vs. closed field lines, for example) plays a very important role and may be such that, at times, parts of the corona become temporarily radiatively efficient. The hard X-ray flaring event recently detected by Chandra from Sgr A^* (Baganoff et al. 2001) might be an example of such an occasional coronal re-brightening. At the very low accretion rates inferred for this source, we predict the X-rays to be produced by inverse Compton scattering on synchrotron photons (see section 4), as indeed is suggested from the observations (Baganoff et al. 2001). We note however that the interpretation of the galactic centre as a coronal-outflow-dominated disc is not straightforward. Firstly, the extreme faintness of the source requires an accretion rate in the inner part of the flow extremely low compared to that expected from simulations of the multiple wind sources in the inner ~ 0.06 pc (Coker & Melia 1997) and values of η very close to unity. The former requirement of small $\dot{m}$ (also consistent with the polarization characteristics of SGR A^* at millimeter and submillimeter wavelength, see Agol, 2000, Quataert & Gruzinov, 2000b, and Melia & Falcke, 2001, and references therein) may be due to the feedback from the coronal outflow suppressing the accretion rate at larger radii. However, a second problem faced by our model is the over-prediction of the IR emission that should be produced by the infalling wind impacting on the outermost part of the cold disc (Falcke & Melia 1997; Coker, Melia & Falcke 1999). Apart from the possibility that the above mentioned interaction between outflowing and inflowing gas modify substantially the kinematics of the accretion flow, it is possible that in fact the inflowing matter forms a Keplerian disc only relatively close to the central black hole (Melia, Liu & Coker 2001). Such a small disc would still be dominated by a strong outflowing corona as the one we have presented here, with enough power to produce a jet responsible for the radio emission, as in the models of Falcke & Markoff (2000). Clearly, the challenge posed by the wealth of data from the galactic centre to any accretion model requires more detailed work than the one we have outlined here, which would be beyond the aim of this paper.

An important related issue concerns the fate of the cold thin disc when we extrapolate our solution to extremely low accretion rates. As we have discussed already in section 2, very strong outflowing coronae might have relevant effects on the nature of the angular momentum transport in the flow, and therefore on the disc structure. The constant improvement of numerical simulations in recent years, make us believe that the relevance of the solutions we describe here will be assessed in the near future. At present, numerical simulations of MRI unstable accretion discs with sufficient extension in the vertical direction seem able to generate only modestly powerful coronae ($f \sim 20$ per cent, Miller & Stone 2000). The inclusion of radiation in the MHD codes (Turner & Stone 2001), and a detailed study of the dependence of the solution on the accretion rate will prove crucial to settle the issue.

We have discussed separately two different outflow–producing mechanisms (MHD and thermal). Their relative importance crucially depends on the uncertain nature of the magnetic energy dissipation in the corona and of particle heating. However, as we have discussed in section 3, according to our current theoretical understanding of these issues, at least one of the two mechanisms must be at work in the physical conditions we envisage for a powerful corona. Clearly both may be at work at the same time, and influence each other. For example, when the conditions for MHD jet launching are favourable, and a larger number of open field lines thread the inner part of the corona, then, depending on the inclination angle of the field lines, a larger mass can be channeled into the outflow, reducing the coronal density and therefore the Coulomb cooling of the hot ions and enhancing the thermally driven jet power.

Radiatively inefficient flows are also expected at high accretion rate, when the flow is geometrically and optically thick and the radiation produced is trapped by the flow and advected with it (Katz 1977; Begelman & Meier 1982; Beloborodov 1998). When the accretion rate is of the order of the Eddington one, the disc thickness stays moderate and a vertically integrated approximation may be retained (slim discs, Abramowicz et al. 1988). In the limit $\dot{m} \gg 1$, though, the behaviour of the disc still remains an open issue. The problem is inherently 2D (Igumenshchev & Abramowicz 2000), and the simultaneous roles of convection, advection and outflows have to be assessed in order to properly model the expected SED. If a strong outflow was produced, through a magnetic corona analogous to the one we envisaged for low accretion rates, or through any other radiative or hydrodynamical mechanism, the total amount of outflowing kinetic power could exceed the Eddington limit (and be much larger than the radiated one; Meier 1982). We just remark here that observationally, black holes believed to accrete at super-Eddington rates bear many similarities with their low-accretion rate counterparts. This class of object may include GBHC in the very high state (see Done 2001 and references therein), whose X-ray spectra show a power-law component usually as strong as the quasi-thermal one; the microquasar GRS 1915+105 (Belloni, Migliari, & Fender 2000); broad line radio galaxies (BLRG, see

e. g. Zdziarski 1999 and reference therein); powerful blazars (see in particular Ghisellini & Celotti 2001a, where it is shown how the relativistic jet in these sources may dominate the total output power); and possibly, narrow line Seyfert 1s (NLS1, Mineshige et al. 2000). Most likely, newly formed supermassive black holes at the epoch of galaxy formation and quasars in their early evolutionary stages might be fed at super-Eddington rates, therefore a better understanding of accretion flows and jet/outflow production in these sources will shed light on the processes of galaxy formation and evolution (Fabian 1999).

6 CONCLUSIONS

We studied the energetics and power output of magnetic structured coronae generated by geometrically thin, optically thick accretion discs at low accretion rates. We have assumed that angular momentum transport in the disc is due to magnetic turbulent stresses. The magnetic field, amplified by MRI, saturates because of the strong buoyancy of flux tubes. Then, we have shown that both magnetic energy density and effective viscous stresses inside the disc are proportional to the geometric mean of the total (gas plus radiation) and gas pressure. In this case, the relative strength of the corona increases as the accretion rate decreases. We have then studied the energetics of the corona itself, and assessed the relevance of the various mechanisms that can generate powerful outflows from the coronal region. We have shown that, depending on the actual coronal field geometry, MHD launching mechanisms can indeed produce strong outflows, whose power is of the order of the radiated one. Moreover, if the corona is two-temperature, a thermally driven outflow can easily dominate the power output from the source at low accretion rates. Thus, if the jet/outflow do not radiate efficiently, the whole system can be very under-luminous compared to a standard radiatively efficient accretion disc. The model we have presented might be relevant to the accretion mode of low-luminosity black holes, either GBHC in the *low/hard* state or LLAGN observed in the nuclei of nearby galaxies.

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